Prediction of $p\bar{\Omega}$ States and Femtoscopic Study

Authors: Yan, Dr. Ye, Huang, Dr. Qi, Dr. Qian Wu, Huang, Prof. Hongxia, Ping, Prof. Jialun, Huang, Prof. Hongxia

Date: 2025-09-12T11:19:48+00:00

Abstract

Inspired by recent research on the $p\Omega$ and $p\bar{\Lambda}$ systems, we investigate the $p\bar{\Omega}$ systems within the framework of the quark delocalization color screening model. Our result indicates that the nucleon-$\bar{\Omega}$ interaction is slightly stronger than the nucleon-Ω interaction, implying a higher likelihood for the $p\bar{\Omega}$ system to form bound states. Dynamic calculations show that the $p\bar{\Omega}$ systems with $J^P = 1^-$ and 2^- form bound states, whose binding energies are deeper than that of the $p\Omega$ system with $J^P = 2^+$. The scattering phase shifts and extracted scattering parameters also support the existence of $p\bar{\Omega}$ bound states. Additionally, we discuss the behavior of the femtoscopic correlation function for the $p\bar{\Omega}$ pairs for the first time. Building on the recent experimental progress on the $p\Omega$ correlation function, future femtoscopic investigations of the $p\bar{\Omega}$ system in heavy-ion collisions will be particularly valuable for constraining baryon-antibaryon interactions.

Full Text

Prediction of $p\bar{\Omega}$ States and Femtoscopic Study

Ye Yan[1,2], **Qi Huang**[1], **Qian Wu**[3,4], **Hong-Xia Huang**[1,†], and **Jia-Lun Ping**[1]

[1]Department of Physics, Nanjing Normal University, Nanjing 210023, China
[2]Changzhou College of Information Technology, Changzhou 213164, China
[3]School of Physics, Nanjing University, Nanjing 210000, China
[4]Institute of Modern Physics, Chinese Academy of Sciences, Lanzhou 730000, China

Inspired by recent research on the $p\Omega$ and $p\bar{\Lambda}$ systems, we investigate the $p\bar{\Omega}$ system within the framework of the quark delocalization color screening model. Our results indicate that the nucleon-$\bar{\Omega}$ interaction is slightly stronger than the nucleon-Ω interaction, implying a higher likelihood for the $p\bar{\Omega}$ system to form bound states. Dynamic calculations show that the $p\bar{\Omega}$ systems with $J^P = 1^-$ and 2^- form bound states whose binding energies are deeper than that of the $p\Omega$ system with $J^P = 2^+$. The scattering phase shifts and extracted scattering parameters also support the existence of $p\bar{\Omega}$ bound states. Additionally, we discuss for the first time the behavior of the femtoscopic correlation function for $p\bar{\Omega}$ pairs. Building on recent experimental progress on the $p\Omega$ correlation function, future femtoscopic investigations of the $p\bar{\Omega}$ system in heavy-ion collisions will be particularly valuable for constraining baryon-antibaryon interactions.

Keywords: $p\bar{\Omega}$ systems; femtoscopic correlation function; bound states; hadron-hadron interaction; scattering phase shifts

Introduction

The study of baryon-antibaryon bound states dates back to the proposal by Fermi and Yang [1] to form the pion from a nucleon-antinucleon pair. In the traditional one-boson-exchange theory of nucleon-nucleon interactions, it is shown that the nucleon-antinucleon system is more attractive than the nucleon-nucleon system due to strong ω-exchange [2]. Therefore, possible bound states and resonances of nucleon-antinucleon systems have been proposed for many years, with an extensive and comprehensive review of possible $N\bar{N}$ bound states provided in Ref. [3].

More recently, the BESIII Collaboration reported the observation of a new $X(1880)$ state in the line shape of the $3(\pi^+\pi^-)$ invariant mass spectrum [4], which is considered evidence for the existence of a proton-antiproton bound state. This has sparked many theoretical works studying the $p\bar{p}$ system and the properties of $X(1880)$ [5-11]. In addition to the possible proton-antiproton state, there has also been great progress in recent work related to $p\bar{\Lambda}$ [12]. A narrow structure in the $p\bar{\Lambda}$ system near the mass threshold, named $X(2085)$, is observed in the process $e^+e^- \to pK^-\bar{\Lambda}$ with a statistical significance exceeding 20σ. Its spin and parity are slightly favored to be $J^P = 1^+$ through an amplitude analysis, with further theoretical results and discussions available in Refs. [13-15]. Building on this significant progress in studies of nucleon-antinucleon and nucleon-$\bar{\Lambda}$ systems, it is natural to explore whether bound states or resonance states can be formed between a nucleon and other hyperons, or between a nucleon and other antihyperons.

In recent years, progress in understanding the strange dibaryon $p\Omega$ has renewed interest in dibaryon systems. The STAR Collaboration measured the $p\Omega$ correlation functions in Au+Au collisions at the Relativistic Heavy-Ion Collider (RHIC) [16] and reported a positive scattering length for the $p\Omega$ interaction, which supports the hypothesis of a $p\Omega$ bound state. In addition, the ALICE Collaboration reported measurements of the p-Ω correlation in pp collisions at $\sqrt{s} = 13$ TeV at the Large Hadron Collider (LHC) [17]. Beyond the $p\Omega$ system, femtoscopic techniques and correlation function studies have made significant progress both experimentally [18-24] and theoretically [25-57].

The $S = -3$, $I = 1/2$, $J = 2$ $N\Omega$ state was first predicted by J. T. Goldman et al. as a narrow resonance in a relativistic quark model [58]. M. Oka also proposed the existence of a quasi-bound state with $I(J^P) = 1/2(2^+)$ using a constituent quark model [59]. A lattice QCD study by the HAL QCD Collaboration reported that the $p\Omega$ state is a bound state at a pion mass of 875 MeV [60], with the bound nature later confirmed with nearly physical quark masses ($m_\pi \simeq 146$ MeV and $m_K \simeq 525$ MeV) [61]. Using interactions obtained from $(2+1)$-flavor lattice QCD simulations, K. Morita et al. studied the two-pair momentum correlation functions of the $p\Omega$ state in relativistic heavy-ion collisions to further investigate the existence of a $p\Omega$ bound state [62, 63]. This state has also been confirmed as a bound state in the frameworks of the chromomagnetic model [64], QCD sum rules [65], and other quark models [66–68]. Additionally, studies on the production of $p\Omega$ and $NN\Omega$ systems can be found in Refs. [69–71].

By analogy to nucleon-nucleon and nucleon-antinucleon systems, one may expect attractive interactions in both the $p\Omega$ and $p\bar{\Omega}$ channels. If the $p\Omega$ state can be confirmed through further experimental measurements, we hope to observe an even stronger signal for the $p\bar{\Omega}$ state in nucleon-antinucleon experiments. In contrast, the $p\bar{\Omega}$ state cannot annihilate into the vacuum, as the nucleon consists of three $u(d)$ quarks and $\bar{\Omega}$ consists of three $\bar{s}$ quarks. In this context, the $p\bar{\Omega}$ state is expected to be relatively stable and may serve as an ideal system for studying baryon-antibaryon interactions. The copious production of antibaryons in high-energy colliders offers excellent opportunities to study this type of spectrum. Clearly, the theoretical study of the $p\bar{\Omega}$ system is both interesting and necessary, as it can provide valuable insights for experimental searches of baryon-antibaryon bound states.

In our previous work, we studied the $p\Omega$ interactions and correlation functions based on the quark delocalization color screening model (QDCSM) [72]. According to our calculations, the depletion of the $p\Omega$ correlation functions caused by the $J^P = 2^+$ bound state, which is not observed in the ALICE Collaboration' s measurements [17], can be explained by the contribution of the attractive $J^P = 1^+$ component in spin-averaging. The QDCSM is a constituent quark model [73, 74] that introduces two key ingredients: first, quark delocalization, which accounts for orbital excitation by allowing quarks to delocalize from one cluster to another; second, the color screening factor, which modifies the confinement interaction between quarks in different cluster orbits. In the study of nucleon-nucleon and nucleon-hyperon interactions and the properties of the deuteron, the mechanism of quark delocalization and color screening plays a crucial role in generating intermediate-range attraction [75–77]. This model has also been used to investigate various dibaryon candidates, such as d^* [78], $p\Omega$ [67, 72], and others [79–83]. It has been extended to study baryon-antibaryon systems, including $p\bar{p}$ and $p\bar{\Lambda}$ [84, 85]. Extending it to the $p\bar{\Omega}$ system is a natural progression. Therefore, we continue to investigate the $p\bar{\Omega}$ system within the framework of the QDCSM, studying it from three aspects: energy spectrum, scattering processes, and correlation functions.

This paper is organized as follows. A brief introduction of the QDCSM is given in the next section. The correlation function and the inverse scattering method are introduced in Sec. II B and Sec. II C, respectively. Sec. III is devoted to numerical results and discussions. The summary is presented in the last section.

Theoretical Formalism

A. Quark Delocalization Color Screening Model

The details of the QDCSM employed in the present work can be found in Refs. [73-77]. Here, we present the salient features of the model. The model Hamiltonian is given by:

$$H = \sum_{i=1}^{6}\left(m_i + \frac{p_i^2}{2m_i}\right) - T_{\text{c.m.}} + \sum_{j>i=1}^{3} V_{qq}(r_{ij}) + \sum_{j>i=4}^{6} V_{\bar{q}\bar{q}}(r_{ij}) + \sum_{\substack{i=1,2,3\\ j=4,5,6}} V_{q\bar{q}}(r_{ij}),$$

including color confinement (V_{CON}), perturbative one-gluon exchange interaction (V_{OGE}), and dynamical chiral symmetry breaking (V_χ):

$$V(r_{ij}) = V_{\text{CON}}(r_{ij}) + V_{\text{OGE}}(r_{ij}) + V_\chi(r_{ij}).$$

A phenomenological color screening confinement potential (V_{CON}) is used as:

$$V_{\text{CON}}(r_{ij}) = -a_c \lambda_i^c \cdot \lambda_j^c \left[f(r_{ij}) + V_0\right],$$

where

$$f(r_{ij}) = \begin{cases} r_{ij}^2 & \text{if } i,j \text{ occur in the same cluster} \\ \frac{1-e^{-\mu_{q_i q_j} r_{ij}^2}}{\mu_{q_i q_j}} & \text{if } i,j \text{ occur in different clusters} \end{cases}$$

Here, a_c, V_0, and $\mu_{q_i q_j}$ are model parameters, and λ^c stands for the SU(3) color Gell-Mann matrices. The color screening parameter $\mu_{q_i q_j}$ is determined by fitting the deuteron properties, nucleon-nucleon scattering phase shifts, and hyperon-nucleon scattering phase shifts, respectively, with $\mu_{qq} = 0.45$, $\mu_{qs} = 0.19$, and $\mu_{ss} = 0.08$ fm^{-2}, satisfying the relation $\mu_{qs}^2 = \mu_{qq}\mu_{ss}$ [86].

The one-gluon exchange potential (V_{OGE}) is written as:

$$V_{\text{OGE}}(r_{ij}) = \alpha_s^{q_i q_j}\left[\lambda_i^c \cdot \lambda_j^c \left(\frac{\pi}{2}\right)\left(\frac{1}{r_{ij}} - \frac{1}{6m_i m_j}\sigma_i \cdot \sigma_j\right)\right],$$

where σ represents the Pauli matrices and α_s is the quark-gluon coupling constant. To cover the wide energy range from light to strange quarks, an effective scale-dependent quark-gluon coupling $\alpha_s(\mu)$ was introduced [87]:

$$\alpha_s(\mu) = \frac{1}{\ln\left(\frac{\mu^2+\mu_0^2}{\Lambda_0^2}\right)}.$$

The dynamical breaking of chiral symmetry results in SU(3) Goldstone boson exchange interactions appearing between constituent light quarks u, d, and s. Hence, the chiral interaction is expressed as $V_\chi(r_{ij}) = V_\pi(r_{ij})+V_K(r_{ij})+V_\eta(r_{ij})$, where:

$$V_\pi(r_{ij}) = \frac{g_{ch}^2}{4\pi}\frac{m_\pi^2}{12m_im_j}\frac{\Lambda_\pi^2}{\Lambda_\pi^2-m_\pi^2}Y(m_\pi r_{ij})(\sigma_i\cdot\sigma_j)\sum_{a=1}^{3}(\lambda_i^a\cdot\lambda_j^a),$$

$$V_K(r_{ij}) = \frac{g_{ch}^2}{4\pi}\frac{m_K^2}{12m_im_j}\frac{\Lambda_K^2}{\Lambda_K^2-m_K^2}Y(m_K r_{ij})(\sigma_i\cdot\sigma_j)\sum_{a=4}^{7}(\lambda_i^a\cdot\lambda_j^a),$$

$$V_\eta(r_{ij}) = \frac{g_{ch}^2}{4\pi}\frac{m_\eta^2}{12m_im_j}\frac{\Lambda_\eta^2}{\Lambda_\eta^2-m_\eta^2}Y(m_\eta r_{ij})(\sigma_i\cdot\sigma_j)\left[\cos\theta_p(\lambda_i^8\cdot\lambda_j^8)-\sin\theta_p(\lambda_i^0\cdot\lambda_j^0)\right],$$

where $Y(x) = e^{-x}/x$ is the standard Yukawa function. The physical η meson is considered by introducing the angle θ_p instead of the octet one. The λ^a are the SU(3) flavor Gell-Mann matrices. The values of m_π, m_K, and m_η are the masses of the SU(3) Goldstone bosons, which adopt the experimental values [88]. The chiral coupling constant g_{ch} is determined from the πNN coupling constant through:

$$\frac{g_{ch}^2}{4\pi} = \left(\frac{3}{5}\right)^2\frac{g_{\pi NN}^2}{4\pi}.$$

Assuming that flavor SU(3) is an exact symmetry broken only by the different mass of the strange quark, the other symbols in the above expressions have their usual meanings.

Regarding $V_{\bar{q}\bar{q}}(r_{ij})$ and $V_{q\bar{q}}(r_{ij})$ in Eq. (1), which represent the antiquark-antiquark ($\bar{q}\bar{q}$) and quark-antiquark ($q\bar{q}$) interactions, the forms of $V_{\bar{q}\bar{q}}$ and $V_{q\bar{q}}$ can be derived by replacing λ_i in Eqs. (3) and (4) with $-\lambda_i^*$ for the antiquark. Notably, there is no annihilation between quark and antiquark because the $p\bar{\Omega}$ state cannot annihilate to the vacuum due to the different quark flavor contents of N and $\bar{\Omega}$. All parameters are taken from our previous work on the $p\Omega$ systems [67], fixed by masses of the ground-state baryons.

In addition, quark delocalization was introduced to enlarge the model variational space to account for mutual distortion or internal excitations of nucleons during interaction. It is realized by specifying the single-particle orbital wave function of the QDCSM as a linear combination of left and right Gaussians:

$$\psi_\alpha(S_i,\epsilon) = \frac{\phi_\alpha(S_i) + \epsilon\phi_\alpha(-S_i)}{N(\epsilon)},$$

$$\psi_\beta(-S_i,\epsilon) = \frac{\phi_\beta(-S_i) + \epsilon\phi_\beta(S_i)}{N(\epsilon)},$$

$$N(S_i,\epsilon) = \sqrt{1+\epsilon^2+2\epsilon e^{-{S_i}^2/4b^2}}.$$

The mixing parameter ϵ is not an adjusted parameter but is determined variationally by the dynamics of the multiquark system itself. In this way, the multiquark system chooses its favorable configuration in the interacting process. This mechanism has been used to explain the crossover transition between the hadron phase and quark-gluon plasma phase [89].

B. Two-Particle Correlation Function

Experimentally, the correlation function $C(k)$ can be measured based on:

$$C(k) = \xi(k)\frac{N_{\text{same}}(k)}{N_{\text{mixed}}(k)},$$

where $N_{\text{same}}(k)$ and $N_{\text{mixed}}(k)$ represent the k distributions of hadron-hadron pairs produced in the same and in different collisions, respectively, and $\xi(k)$ denotes corrections for experimental effects.

In theoretical studies, the correlation function can be calculated using the Koonin-Pratt (KP) formula [90–92]:

$$C(k) = \frac{\int d^4x_1 d^4x_2 S_1(x_1,p_1)S_2(x_2,p_2)|\Psi(\mathbf{r},k)|^2}{\int d^4x_1 d^4x_2 S_1(x_1,p_1)S_2(x_2,p_2)},$$

where $S_i(x_i,p_i)$ $(i=1,2)$ is the single-particle source function of hadron i with momentum p_i, $k=(m_2p_1-m_1p_2)/(m_1+m_2)$ is the relative momentum in the center-of-mass frame of the pair $(p_1+p_2=0)$, $\mathbf{r}$ is the relative coordinate with time difference correction, and $\Psi(\mathbf{r},k)$ is the relative wave function in the two-body outgoing state with asymptotic relative momentum k.

In the case where we can ignore the time difference of emission and the momentum dependence of the source, we integrate out the center-of-mass coordinate and obtain:

$$C(k) = \int d\mathbf{r} S_{12}(\mathbf{r}) |\Psi(\mathbf{r}, k)|^2,$$

where $S_{12}(\mathbf{r})$ is the normalized pair source function in the relative coordinate, given by:

$$S_{12}(\mathbf{r}) = \frac{1}{(4\pi R^2)^{3/2}} \exp\left(-\frac{r^2}{4R^2}\right),$$

with R being the size parameter of the source. Thus, two important factors of the correlation function are included in Eq. (13): the collision system, related to the source function $S_{12}(\mathbf{r})$, and the two-particle interaction, embedded in the relative wave function $\Psi(\mathbf{r}, k)$.

For a pair of non-identical particles such as $p\bar{\Omega}$, assuming that only the S-wave part of the wave function is modified by the two-particle interaction, $\Psi(\mathbf{r}, k)$ can be written as:

$$\Psi_{p\bar{\Omega}}(\mathbf{r}, k) = e^{i\mathbf{k}\cdot\mathbf{r}} - j_0(kr) + \psi_{p\bar{\Omega}}(\mathbf{r}, k),$$

where the spherical Bessel function $j_0(kr)$ represents the S-wave part of the non-interacting wave function, and $\psi_{p\bar{\Omega}}$ stands for the scattering wave function affected by the two-particle interaction. Substituting the relative wave function $\Psi_{p\bar{\Omega}}(\mathbf{r}, k)$ into the KP formula yields the correlation function:

$$C_{p\bar{\Omega}}(k) = 1 + \int_0^\infty 4\pi r^2 dr\, S_{12}(r) \left[|\psi_{p\bar{\Omega}}(r, k)|^2 - |j_0(kr)|^2 \right].$$

The scattering wave function $\psi_{p\bar{\Omega}}(r, k)$ can be obtained by solving the Schrödinger equation, and a similar approach has been utilized in femtoscopic correlation analysis tools using the Schrödinger equation [93]:

$$\nabla^2 \psi_{p\bar{\Omega}}(\mathbf{r}, k) + V(\mathbf{r}) \psi_{p\bar{\Omega}}(\mathbf{r}, k) = E \psi_{p\bar{\Omega}}(\mathbf{r}, k),$$

where $\mu = m_p m_\Omega / (m_p + m_\Omega)$ is the reduced mass of the system.

Considering the S-wave case, the wave function can be separated into a radial term $R_k(r)$ and an angular term $Y_0^0(\theta, \phi)$ and expressed as:

$$\psi_{p\bar{\Omega}}(\mathbf{r}, \theta, \phi) = R_k(r) Y_0^0(\theta, \phi).$$

The interaction between a proton and a $\bar{\Omega}$ baryon includes both the strong interaction and the repulsive Coulomb interaction, so the potential can be written as:

$$V(r) = V_{\text{Strong}}(r) + V_{\text{Coulomb}}(r),$$

where $V_{\text{Coulomb}}(r) = +\alpha\hbar c/r$ and α is the fine-structure constant. The method to obtain the strong interaction potential $V_{\text{Strong}}(r)$ will be introduced in the next section.

Once the total interaction potential is determined, the radial Schrödinger equation can be solved:

$$\frac{d^2 u_k(r)}{dr^2} + V(r)u_k(r) = Eu_k(r),$$

where $E = \hbar^2 k^2/(2\mu)$ and $u_k(r) = rR_k(r)$. On this basis, the correlation function $C_{p\bar{\Omega}}(k)$ for given spin-parity quantum numbers can be calculated through Eq. (16). The calculation of correlation functions described above is based on obtaining scattering wave functions by solving the Schrödinger equation in coordinate space [26–28, 30, 32, 33, 36]. Additionally, scattering wave functions can also be obtained by solving the Lippmann-Schwinger (Bethe-Salpeter) equation in momentum space [29, 31, 38, 39]. Further details on correlation functions for various systems can be found in the references mentioned above.

For the S-wave $p\bar{\Omega}$ dibaryon system, the possible spin-parity quantum numbers are $J^P = 1^-$ and 2^-, respectively. Since the experimentally measured correlation function is spin-averaged, the theoretically obtained correlation function should also consider the average over systems with different quantum numbers:

$$C_{p\bar{\Omega}}(k) = \frac{1}{2}\left[C_{p\bar{\Omega}}^{J=1}(k) + C_{p\bar{\Omega}}^{J=2}(k)\right].$$

C. Gel' fand-Levitan-Marchenko Method for Inverse Scattering Problem

To solve Eq. (20), a two-body interaction potential $V(r)$ is absolutely necessary. The QDCSM is actually a treatment for few-body problems, which means directly extracting a two-body interaction potential $V(r)$ from it is not natural. However, the QDCSM can be employed to investigate scattering processes, from which the desired potential can be obtained since hadronization is fully incorporated in the model.

The approach we adopted to extract the two-body equivalent potential $V(r)$ is the GLM method, which is a very powerful tool in inverse scattering theory [94]. It provides a systematic approach to reconstruct an equivalent potential from the scattering data of a specific process, making it a classical "inverse problem." Thus, this method offers another path to understand the nature of two-body interactions.

The key equation of the GLM method used in this work is the Marchenko equation [95, 96], which can be written for the S-wave case in integral form as:

$$K(r,r') + F(r,r') + \int_r^\infty K(r,s)F(s,r')ds = 0.$$

Here, the kernel function $K(r,r')$ is the solution to be determined, and $F(r,r')$ is the inverse Fourier transform of the reflection coefficient:

$$F(r,r') = \frac{1}{2\pi}\int_{-\infty}^{\infty} e^{ikr}\{1-S(k)\}e^{ikr'}dk + \sum_{i=1}^{n} M_i e^{-\kappa_i r}e^{-\kappa_i r'}.$$

The partial-wave scattering matrix $S(k)$ is given by $S(k) = \exp(2i\delta(k))$, where $\delta(k)$ is the scattering phase shift satisfying $k\cot\delta = -1/a_0 + \frac{1}{2}r_{\text{eff}}k^2$. Here, a_0 and r_{eff} represent the scattering length and effective range, respectively. Additionally, n is the number of bound states, κ_i denotes the wavenumber of the i-th bound state, and M_i is the norming constant. After solving the Marchenko equation and obtaining $K(r,r')$, the potential can be reconstructed as:

$$V(r) = -2\frac{d}{dr}K(r,r).$$

One important point to emphasize is that when bound states exist, this method generally cannot give a fully determined potential but ends up with a set of phase-equivalent potentials [97]. However, if one fixes all the M_i in a unique way, such as by calculating from the Jost solution, the obtained potential will be unique for further calculation [98, 99]. Using this method, preparation for further calculation can be done; for a more comprehensive discussion, one can refer to Refs. [94-102].

Results and Discussion

The S-wave $p\bar{\Omega}$ systems with isospin $I = 1/2$ and spin-parity $J^P = 1^-$ and 2^- are investigated on the basis of the QDCSM. To determine whether any bound state exists, a dynamic calculation is performed as a first step. Here we employ the resonating group method (RGM) to solve the bound-state problem, which is briefly introduced in the Appendix, with more details available in Refs. [103-105]. By expanding the relative motion wave function between two hadrons in the RGM equation by Gaussians, the integro-differential equation of RGM can be reduced to an algebraic equation—the generalized eigen-equation. The energy of the system can then be obtained by solving this eigen-equation. The binding energies of the $p\bar{\Omega}$ systems with $J^P = 1^-$ and $J^P = 2^-$, denoted as E_B, are listed in Table 1 .

Here, $E^{\text{th}}_{\text{Theo}}$ represents the theoretical threshold, and $E^{\text{cal}}_{\text{Theo}}$ represents the eigenvalue of the corresponding system. The calculation for the $p\bar{\Omega}$ systems does not involve channel coupling because we limit our study to color-singlet sub-clusters consisting of three u/d quarks and three $\bar{s}$ quarks.

From Table 1, we find that both the $J^P = 1^-$ and $J^P = 2^-$ $p\bar{\Omega}$ systems form bound states with binding energies of about 10 MeV and 9 MeV, respectively. By contrast, in our previous work on the $p\Omega$ systems, the single-channel calculation shows that neither the $J^P = 1^+$ nor $J^P = 2^+$ $p\Omega$ is bound. After channel coupling, only the $p\Omega$ with $J^P = 2^+$ forms a bound state with binding energy of about 6 MeV. These numerical results indicate that it is more likely for the $p\bar{\Omega}$ system than the $p\Omega$ system to form bound states in our calculations. Therefore, considering that attractive $p\Omega$ interaction is implied in experimental measurements of $p\Omega$ correlation functions [17], we look forward to experimental progress on $p\bar{\Omega}$ correlation functions in the future.

To further study the interaction between nucleon and $\bar{\Omega}$, we calculated the scattering phase shifts of the $p\bar{\Omega}$ systems. The calculation is based on the well-developed Kohn-Hulthen-Kato (KHK) variational method, with details of this method available in Refs. [82, 104]. The low-energy scattering phase shifts of the $p\bar{\Omega}$ systems are shown in Fig. 1 [Figure 1: see original paper]. For both $J^P = 1^-$ and $J^P = 2^-$ $p\bar{\Omega}$ systems, the scattering phase shifts approach 180° when $E_{\text{c.m.}} \to 0$ MeV and rapidly decrease when $E_{\text{c.m.}}$ increases, indicating the existence of bound states in these systems. This conclusion is consistent with the bound-state calculation discussed earlier.

We then extract the scattering length a_0 and effective range r_{eff} of the $p\bar{\Omega}$ systems from the low-energy phase shifts using the expansion:

$$k \cot \delta = -\frac{1}{a_0} + \frac{1}{2} r_{\text{eff}} k^2 + O(k^4),$$

where k is the momentum of relative motion with $k = \sqrt{2\mu E_{\text{c.m.}}}$, μ is the reduced mass of the two baryons, and $E_{\text{c.m.}}$ is the incident energy; δ is the low-energy scattering phase shift. The binding energy E'_B can be calculated according to:

$$E'_B = \frac{\hbar^2 \alpha^2}{2\mu},$$

where α is the wave number obtained from the relation [106]:

$$\alpha = \frac{1 - r_{\text{eff}}/a_0 + \sqrt{(1 - r_{\text{eff}}/a_0)^2 + (2r_{\text{eff}}/a_0)^2}}{2r_{\text{eff}}}.$$

Note that this is another way to calculate the binding energy, hence it is labeled E'_B. The scattering parameters of the $p\bar{\Omega}$ systems, along with the binding energies obtained using these parameters, are listed in Table 2 .

Our results show that the scattering lengths are positive for the $p\bar{\Omega}$ systems with $J^P = 1^-$ and 2^-, which also confirms the existence of bound states. Furthermore, the binding energies of the two systems obtained by Eq. (26) are broadly consistent with the numerical results shown in Table 1, which were obtained by dynamic calculation. Additionally, in the method of obtaining binding energies using scattering parameters, the binding energies of the two $p\bar{\Omega}$ systems are also deeper than that of the $p\Omega$ system with $J^P = 2^+$.

Moreover, by solving the inverse scattering problem, we can further study the behavior of the $p\bar{\Omega}$ correlation functions based on the $p\bar{\Omega}$ scattering process and the KP formula in Eq. (13). Before that, we examine the general properties of the $p\bar{\Omega}$ correlation functions through an effective square well potential model. The correlation functions corresponding to different square well potentials are presented in Fig. 2 [Figure 2: see original paper]. The solid red lines represent correlation functions influenced only by the repulsive Coulomb interaction. The dashed blue lines represent correlation functions influenced only by square well potentials. We introduce three square well potentials with width $r_0 = 2$ fm: weak attraction in panel (a) corresponding to $V_0 = -10$ MeV, moderate attraction in panel (b) corresponding to $V_0 = -28$ MeV, and relatively strong attraction in panel (c) corresponding to $V_0 = -40$ MeV. The dotted black lines represent correlation functions influenced by both Coulomb interaction and square well potentials through Eq. (19). The source size parameter R in Eq. (14) is taken from our previous work on $p\Omega$ correlation functions.

In Fig. 2, panels (a) and (b), the correlation function affected only by the square potential is above unity in the low-energy region due to the attractive interaction. The difference is that weak attraction is insufficient to form a bound state, so the correlation function remains above unity, while moderate attraction forms a shallow bound state. The existence of the bound state leads to depletion of the correlation function, resulting in a portion below unity. After accounting for both Coulomb and square well potentials, which dominate the low-energy region ($0 < k < 25$ MeV), the correlation function forms a peak-like structure. In panel (c), the correlation function remains below unity for relatively strong attraction. A discussion of this phenomenon can be found in Refs. [39, 72]. After considering the Coulomb interaction, the correlation function in the low-energy region nearly coincides with the result obtained by considering only the Coulomb interaction. As the relative momentum k increases, the correlation function closely matches the result obtained by considering only the relatively strong attraction.

After replacing the square well potentials with effective potentials obtained by solving the inverse scattering problem using the GLM method (briefly introduced in Sec. II C), we can study the correlation function of the $p\bar{\Omega}$ systems. Both effective potentials for the $p\bar{\Omega}$ systems with $J^P = 1^-$ and 2^- are obtained. The total $p\bar{\Omega}$ correlation function is the superposition of correlation functions corresponding to the two quantum numbers according to Eq. (21). Since the $p\bar{\Omega}$ system forms bound states for both quantum numbers and the interactions are similar, we omit comparison of the correlation functions for the two quantum numbers here. The total correlation functions calculated for different values of source size parameter R are shown in Fig. 3 [Figure 3: see original paper]. Recent studies on emission source properties can be found in Refs. [107–111].

In Fig. 3, panels (a)-(e), the dashed gray lines and dotted orange lines represent the $p\bar{\Omega}$ correlation functions considering only the Coulomb interaction and both Coulomb and strong interactions, respectively. In panel (f), the gray band stands for correlation functions influenced only by the Coulomb interaction with size parameter R ranging from 1.0 to 2.5 fm, while the other lines summarize the correlation functions shown in panels (a)-(e). According to our results, changes in the size parameter R can greatly influence the $p\bar{\Omega}$ correlation functions. An obvious feature is that as R increases, the peak-like structure caused by the different dominant regions of Coulomb and $p\bar{\Omega}$ strong interactions gradually becomes less obvious and eventually disappears. Since two bound states are obtained in our calculation, it is very important to verify this conclusion in the correlation functions. It can be seen that as the correlation function influenced only by the Coulomb interaction gradually approaches unity, the depletion caused by the bound states leads to the correlation function being below that of the Coulomb-only case.

In recent years, experimental data on correlation functions have increased rapidly [18–24], providing unprecedented insights into hadron-hadron interactions across various systems. Together with femtoscopic techniques, these studies open new possibilities for extracting low-energy scattering parameters that are difficult to access otherwise. On the theoretical side, continuous progress in lattice QCD, effective field theory, and quark-model-based approaches has greatly enriched our understanding and offered valuable guidance for interpreting experimental observations [25–57]. The synergy between experimental measurements and theoretical developments will not only deepen our knowledge of the strong interaction but also pave the way for future explorations of exotic hadronic states and hypernuclear physics. In this context, it is well established that hyperon-nucleon interactions provide the basis for hypernuclei formation [112–116], while antihyperon-antinucleon interactions give rise to anti-hypernuclei [117, 118]. An open question is whether antihyperon-nucleon interactions could also lead to novel nuclear systems. As a continuation of the present study, we will further explore such systems based on the antihyperon-nucleon interactions obtained in this work.

Summary

In this work, we investigate the S-wave $p\bar{\Omega}$ systems with isospin $I = 1/2$ and spin-parity $J^P = 1^-$ and 2^- in the framework of the QDCSM. The results show that the $p\bar{\Omega}$ systems with both $J^P = 1^-$ and 2^- form bound states, and the attraction between nucleon and $\bar{\Omega}$ is slightly larger than that between nucleon and Ω, suggesting that the $p\bar{\Omega}$ system has a higher likelihood of forming bound states than the $p\Omega$ system. Calculations of low-energy scattering phase shifts and scattering parameters for the $p\bar{\Omega}$ systems also support the existence of $p\bar{\Omega}$ bound states with $J^P = 1^-$ and 2^-. Furthermore, since the nucleon is composed of three light $u(d)$ quarks and $\bar{\Omega}$ of three strange quarks $\bar{s}$, the $p\bar{\Omega}$ state cannot annihilate to the vacuum. In this context, the $p\bar{\Omega}$ state is special and can provide useful information for experimental searches of baryon-antibaryon bound states.

Using the GLM method, we solve the inverse scattering problem and obtain effective $p\bar{\Omega}$ potentials. On this basis, the $p\bar{\Omega}$ correlation functions are calculated, taking into account both Coulomb interaction and spin averaging. We present correlation functions corresponding to different source size parameters R, which can be used for future comparison with experimental measurements.

Understanding hadron-hadron interactions is one of the important issues in hadron physics. The study of baryon-antibaryon interactions in this work provides an effective way to test this mechanism. The femtoscopic correlation function has become an important tool for exploring hadron-hadron interactions, and further theoretical and experimental investigations are essential for a deeper understanding of such baryon-antibaryon systems.

References

[1] E. Fermi and C. N. Yang, Phys. Rev. 76, 1739 (1949).
[2] K. Erkelenz, Phys. Rept. 13, 191 (1974).
[3] J. M. Richard, Nucl. Phys. B Proc. Suppl. 86, 361 (2000).
[4] M. Ablikim et al. [BESIII], Phys. Rev. Lett. 132, 151901 (2024).
[5] S. G. Salnikov and A. I. Milstein, Nucl. Phys. B 1002, 116539 (2024).
[6] Y. Xiao, J. X. Lu and L. S. Geng, Phys. Rev. C 110, 6 (2024).
[7] M. Karliner and J. L. Rosner, Phys. Rev. D 110, 094058 (2024).
[8] P. Y. Niu, Z. Y. Zhang, Y. Y. Li, Q. Wang and Q. Zhao, Phys. Rev. D 110, 094020 (2024).
[9] P. G. Ortega, D. R. Entem, F. Fernandez and J. Segovia, Phys. Lett. B 862, 139281 (2025).
[10] Z. S. Jia, Z. H. Zhang, F. K. Guo and G. Li, Phys. Rev. D 111, 054014 (2025).
[11] Q. H. Yang, L. Y. Dai and U. G. Meißner, [arXiv:2412.07599 [hep-ph]].
[12] M. Ablikim et al. [BESIII], Phys. Rev. Lett. 131, 151901 (2023).
[13] T. G. Li, S. C. Zhang, G. Y. Wang and Q. F. Lü, Phys. Rev. D 110, 114020 (2024).
[14] J. Haidenbauer and U. G. Meißner, Eur. Phys. J. A 60, 119 (2024).
[15] X. H. Zhang, S. Q. Zhang and C. F. Qiao, Eur. Phys. J. C 85, 693 (2025).
[16] J. Adam et al. [STAR], Phys. Lett. B 790, 490 (2019).
[17] A. Collaboration et al. [ALICE], Nature 588, 232 (2020) [erratum: Nature 590, E13 (2021)].
[18] L. Adamczyk et al. [STAR], Nature 527, 345 (2015).
[19] S. Acharya et al. [ALICE], Phys. Rev. Lett. 124, 092301 (2020).
[20] L. Fabbietti, V. Mantovani Sarti and O. Vazquez Doce, Ann. Rev. Nucl. Part. Sci. 71, 377 (2021).

[21] S. Acharya et al. [ALICE], Phys. Rev. Lett. 127, 172301 (2021).
[22] S. Acharya et al. [ALICE], Phys. Lett. B 833, 137272 (2022).
[23] S. Acharya et al. [ALICE], Phys. Rev. D 106, 052010 (2022).
[24] B. E. Aboona et al. [STAR], Phys. Lett. B 864, 139412 (2025).
[25] A. Ohnishi, Y. Hirata, Y. Nara, S. Shinmura and Y. Akaishi, Nucl. Phys. A 670, 297 (2000).
[26] K. Morita, T. Furumoto and A. Ohnishi, Phys. Rev. C 91, 024916 (2015).
[27] A. Ohnishi, K. Morita, K. Miyahara and T. Hyodo, Nucl. Phys. A 954, 294 (2016).
[28] T. Hatsuda, K. Morita, A. Ohnishi and K. Sasaki, Nucl. Phys. A 967, 856 (2017).
[29] J. Haidenbauer, Nucl. Phys. A 981, 1 (2019).
[30] Y. Kamiya, T. Hyodo, K. Morita, A. Ohnishi and W. Weise, Phys. Rev. Lett. 124, 132501 (2020).
[31] J. Haidenbauer, G. Krein and T. C. Peixoto, Eur. Phys. J. A 56, 184 (2020).
[32] A. Ohnishi, Y. Kamiya, K. Sasaki, T. Fukui, T. Hatsuda, T. Hyodo, K. Morita and K. Ogata, Few Body Syst. 62, 42 (2021).
[33] K. Ogata, T. Fukui, Y. Kamiya and A. Ohnishi, Phys. Rev. C 103, 065205 (2021).
[34] S. Mrówczyński and P. Słoń, Phys. Rev. C 104, 024909 (2021).
[35] Ł. K. Graczykowski and M. A. Janik, Phys. Rev. C 104, 054909 (2021).
[36] Y. Kamiya, K. Sasaki, T. Fukui, T. Hyodo, K. Morita, K. Ogata, A. Ohnishi and T. Hatsuda, Phys. Rev. C 105, 014915 (2022).
[37] J. Haidenbauer and U. G. Meißner, Phys. Lett. B 829, 137074 (2022).
[38] Z. W. Liu, K. W. Li and L. S. Geng, Chin. Phys. C 47, 024108 (2023).
[39] Z. W. Liu, J. X. Lu and L. S. Geng, Phys. Rev. D 107, 074019 (2023).
[40] Z. W. Liu, J. X. Lu, M. Z. Liu and L. S. Geng, Phys. Rev. D 108, L031503 (2023).
[41] R. Molina, Z. W. Liu, L. S. Geng and E. Oset, Eur. Phys. J. C 84, 328 (2024).
[42] I. Vidana, A. Feijoo, M. Albaladejo, J. Nieves and E. Oset, Phys. Lett. B 846, 138201 (2023).
[43] V. M. Sarti, A. Feijoo, I. Vidaña, A. Ramos, F. Giacosa, T. Hyodo and Y. Kamiya, Phys. Rev. D 110, L011505 (2024).
[44] R. Molina, C. W. Xiao, W. H. Liang and E. Oset, Phys. Rev. D 109, 054002 (2024).
[45] M. Albaladejo, A. Feijoo, I. Vidaña, J. Nieves and E. Oset, Eur. Phys. J. A 61, 187 (2025).
[46] A. Feijoo, L. R. Dai, L. M. Abreu and E. Oset, Phys. Rev. D 109, 016014 (2024).
[47] H. P. Li, J. Y. Yi, C. W. Xiao, D. L. Yao, W. H. Liang and E. Oset, Chin. Phys. C 48, 053107 (2024).
[48] A. Feijoo, M. Korwieser and L. Fabbietti, Phys. Rev. D 111, 014009 (2025).
[49] N. Ikeno, G. Toledo and E. Oset, Phys. Lett. B 847, 138281 (2023).
[50] M. Albaladejo, J. Nieves and E. Ruiz-Arriola, Phys. Rev. D 108, 014020

(2023).
[51] J. M. Torres-Rincon, À. Ramos and L. Tolos, Phys. Rev. D 108, 096008 (2023).
[52] L. M. Abreu, P. Gubler, K. P. Khemchandani, A. Martinez Torres and A. Hosaka, Phys. Lett. B 860, 139175 (2025).
[53] H. P. Li, C. W. Xiao, W. H. Liang, J. J. Wu, E. Wang and E. Oset, Phys. Rev. D 110, 114018 (2024).
[54] M. Albaladejo, A. Feijoo, J. Nieves, E. Oset and I. Vidaña, Phys. Rev. D 110, 114052 (2024).
[55] F. Etminan, Phys. Rev. C 111, 014912 (2025).
[56] A. Jinno, Y. Kamiya, T. Hyodo and A. Ohnishi, J. Subatomic Part. Cosmol. 1, 100005 (2024).
[57] F. Etminan, Phys. Lett. B 866, 139564 (2025).
[58] J. T. Goldman, K. Maltman, G. J. Stephenson, K. E. Schmidt and F. Wang, Phys. Rev. Lett. 59, 627 (1987).
[59] M. Oka, Phys. Rev. D 38, 298 (1988).
[60] F. Etminan et al. [HAL QCD], Nucl. Phys. A 928, 89 (2014).
[61] T. Iritani et al. [HAL QCD], Phys. Lett. B 792, 284 (2019).
[62] K. Morita, A. Ohnishi, F. Etminan and T. Hatsuda, Phys. Rev. C 94, 031901 (2016).
[63] K. Morita, S. Gongyo, T. Hatsuda, T. Hyodo, Y. Kamiya and A. Ohnishi, Phys. Rev. C 101, 015201 (2020).
[64] B. Silvestre-Brac and J. Leandri, Phys. Rev. D 45, 4221 (1992).
[65] X. H. Chen, Q. N. Wang, W. Chen and H. X. Chen, Phys. Rev. D 103, 094011 (2021).
[66] L. R. Dai, D. Zhang, C. R. Li and L. Tong, Chin. Phys. Lett. 24, 389 (2007).
[67] H. Huang, J. Ping and F. Wang, Phys. Rev. C 92, 065202 (2015).
[68] Q. B. Li and P. N. Shen, Eur. Phys. J. A 8, 417 (2000).
[69] S. Zhang and Y. G. Ma, Phys. Lett. B 811, 135867 (2020).
[70] J. Pu, K. J. Sun, C. W. Ma and L. W. Chen, Phys. Rev. C 110, 024908 (2024).
[71] L. Zhang, S. Zhang and Y. G. Ma, Eur. Phys. J. C 82, 416 (2022).
[72] Y. Yan, Q. Huang, Y. Yang, H. Huang and J. Ping, Sci. China Phys. Mech. Astron. 68, 232012 (2025).
[73] F. Wang, G. h. Wu, L. j. Teng and J. T. Goldman, Phys. Rev. Lett. 69, 2901 (1992).
[74] G. H. Wu, L. J. Teng, J. L. Ping, F. Wang and J. T. Goldman, Phys. Rev. C 53, 1161 (1996).
[75] J. L. Ping, F. Wang and J. T. Goldman, Nucl. Phys. A 657, 95 (1999).
[76] G. h. Wu, J. L. Ping, L. j. Teng, F. Wang and J. T. Goldman, Nucl. Phys. A 673, 279 (2000).
[77] H. R. Pang, J. L. Ping, F. Wang and J. T. Goldman, Phys. Rev. C 65, 014003 (2002).
[78] J. L. Ping, H. X. Huang, H. R. Pang, F. Wang and C. W. Wong, Phys. Rev. C 79, 024001 (2009).

[79] Y. Yan, Y. Wu, X. Hu, H. Huang and J. Ping, Phys. Rev. D 105, 014027 (2022).
[80] Y. Yan, X. Hu, Y. Wu, H. Huang, J. Ping and Y. Yang, Eur. Phys. J. C 83, 524 (2023).
[81] Y. Yan, Y. Wu, H. Huang, J. Ping and X. Zhu, Eur. Phys. J. C 83, 610 (2023).
[82] Y. Yan, X. Hu, H. Huang and J. Ping, Phys. Rev. D 108, 094045 (2023).
[83] Y. Yan, H. Huang, X. Zhu and J. Ping, Phys. Rev. D 109, 034036 (2024).
[84] H. Huang, H. Pang and J. Ping, Mod. Phys. Lett. A 26, 1231 (2011).
[85] H. Huang, J. Ping and F. Wang, Mod. Phys. Lett. A 27, 1250039 (2012).
[86] M. Chen, H. Huang, J. Ping and F. Wang, Phys. Rev. C 83, 015202 (2011).
[87] J. Vijande, F. Fernandez and A. Valcarce, J. Phys. G 31, 481 (2005).
[88] S. Navas et al. [Particle Data Group], Phys. Rev. D 110, 030001 (2024).
[89] M. Xu, M. Yu and L. Liu, Phys. Rev. Lett. 100, 092301 (2008).
[90] S. E. Koonin, Phys. Lett. B 70, 43 (1977).
[91] S. Pratt, T. Csorgo and J. Zimanyi, Phys. Rev. C 42, 2646 (1990).
[92] W. Bauer, C. K. Gelbke and S. Pratt, Ann. Rev. Nucl. Part. Sci. 42, 77 (1992).
[93] D. L. Mihaylov, V. Mantovani Sarti, O. W. Arnold, L. Fabbietti, B. Hohlweger and A. M. Mathis, Eur. Phys. J. C 78, 394 (2018).
[94] K. Chadan and P. C. Sabatier, Inverse Problems in Quantum Scattering Theory, 2nd ed. (Springer, New York, 1989).
[95] V. A. Marchenko, Dokl. Akad. Nauk SSSR, 104, 695 (1955).
[96] Z. S. Agranovich and V. A. Marchenko, The Inverse Problem of the Scattering Theory (Gordon and Breach, New York, 1963).
[97] S. A. Sofianos, A. Papastylianos, H. Fiedeldey and E. O. Alt, Phys. Rev. C 42, R506 (1990).
[98] S. E. Massen, S. A. Sofianos, S. A. Rakityansky and S. Oryu, Nucl. Phys. A 654, 597 (1999).
[99] R. G. Newton, Scattering Theory of Waves and Particles, 2nd ed. (Dover, New York, 2002).
[100] L. Jade, M. Sander and H. V. von Geramb, Lect. Notes Phys. 488, 124 (1997).
[101] N. A. Khokhlov and L. I. Studenikina, Phys. Rev. C 104, 014001 (2021).
[102] N. A. Khokhlov, Phys. Rev. C 107, 044001 (2023).
[103] J. A. Wheeler, Phys. Rev. 52, 1083 (1937).
[104] M. Kamimura, Prog. Theor. Phys. Suppl. 62, 236 (1977).
[105] Y. Yan, Q. Huang, X. Zhu, H. Huang and J. Ping, Phys. Rev. D 110, 014021 (2024).
[106] V. A. Babenko and N. M. Petrov, Phys. Atom. Nucl. 66, 1319 (2003).
[107] S. Acharya et al. [ALICE], Phys. Lett. B 811, 135849 (2020).
[108] S. Acharya et al. [ALICE], Eur. Phys. J. C 85, 198 (2025).
[109] D. Mihaylov and J. González González, Eur. Phys. J. C 83, 590 (2023).
[110] O. Vázquez Doce, D. Mihaylov and L. Fabbietti, Eur. Phys. J. A 61, 53 (2025).
[111] D. F. Wang, M. Y. Chen, Y. G. Ma, Q. Y. Shou, S. Zhang and L. Zheng,

Nucl. Sci. Tech. 36, 154 (2025).
[112] J. Haidenbauer, U. G. Meißner and A. Nogga, Eur. Phys. J. A 56, 91 (2020).
[113] B. Aboona et al. [STAR], Phys. Rev. Lett. 130, 212301 (2023).
[114] J. Chen, X. Dong, Y. G. Ma and Z. Xu, Sci. Bull. 68, 3252-3260 (2023).
[115] S. Acharya et al. [ALICE], Phys. Rev. Lett. 131, 102302 (2023).
[116] Y. G. Ma, Nuc. Sci. Tech. 34, 97 (2023).
[117] B. I. Abelev et al. [STAR], Science 328, 58 (2010).
[118] M. Abdulhamid et al. [STAR], Nature 632, 1026 (2024).

Note: Figure translations are in progress. See original paper for figures.